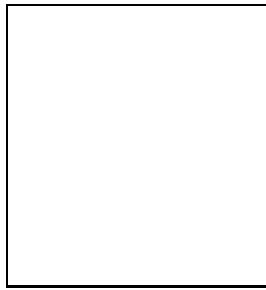


Higgs-less Higgs mechanism: low-energy expansion^aJohannes Hirn[§] and Jan Stern[¶][§]*IFIC, Departament de Física Teòrica, CSIC - Universitat de València
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In this talk, we describe an effective theory for electroweak symmetry breaking without a physical Higgs, based on a symmetry larger than the electroweak gauge group. This symmetry forbids deviations from the Standard Model at the leading order in the appropriate chiral expansion. Indeed, the large symmetry allows for a consistent expansion of the effective theory in powers of momenta and spurions. The latter are automatically present: they define the covariant reduction from the large symmetry to the electroweak group.

1 Introduction

We consider electroweak symmetry breaking (EWSB) without a Higgs boson: the Higgs mechanism removes the three Goldstone bosons (GBs) from the spectrum, while the usual Higgs boson does not exist. We wish to construct a low-energy effective theory (LEET), expanding in powers of momenta (and other naturally-small parameters, the spurions, as we shall soon see). Indeed, for the expansion to be consistent according to the rules of a LEET (such as Chiral Perturbation Theory^{1,2)}, we need to assume the presence of a symmetry S_{nat} controlling the smallness of deviations from the Standard Model (SM), i.e. technical naturalness.

We require the “hidden” symmetry S_{nat} to be sufficiently large as to force the leading $\mathcal{O}(p^2)$ order of the LEET to coincide with the tree-level Higgs-less vertices of the SM in the limit of vanishing fermion masses. The S_{nat} symmetry can be viewed as defining the custodial symmetry (and its extension to the left-handed non-abelian sector) in the presence of non-vanishing gauge coupling g' (g). S_{nat} necessarily contains the SM gauge group $SU(2)_L \times U(1)_Y$, which we call S_{red} . The higher symmetry $S_{\text{nat}} \supset S_{\text{red}}$ can be linearized by adding a set of nine auxiliary gauge

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fields to the original four present in the SM. The additional gauge fields are not physical: they will be eliminated by constraints, implemented via spurions.

Once the constraints are applied, the theory only contains $SU(2)_L \times U(1)_Y$ Yang-Mills fields and chiral fermions coupled to three Goldstone bosons $\Sigma(x)$. The latter disappear from the spectrum, resulting in three of the vector fields acquiring a mass. All physical degrees of freedom are light compared to the scale $\Lambda_w \simeq 4\pi v \simeq 3$ TeV: the only new particles beyond those already known are *light* right-handed neutrinos.

This talk is based on the first half of³. In that paper, we also studied in detail some of the consequences for lepton-number violation (LNV) processes, vertex corrections (which come in before oblique corrections in the present case) as well as the possible contribution of light right-handed neutrinos to dark matter.

2 Higgs-less LEET based only on $SU(2)_L \times U(1)_Y$

2.1 Power-counting

In the case of Higgs-less EWSB, we do not have a renormalizable theory to start with: all operators respecting the symmetries must be included in the effective lagrangian. The infinite number of them should be ordered according to their importance in the low-energy limit $p \rightarrow 0$, i.e. according to their infrared (or chiral) dimension d_{IR} . The effective lagrangian is then expressed as $\mathcal{L}_{eff} = \sum_{d_{IR} \geq 2} \mathcal{L}_{d_{IR}}$ with $\mathcal{L}_{d_{IR}} = \mathcal{O}(p^{d_{IR}})$. A local operator/interaction vertex \mathcal{O} built from GBs, gauge fields and fermions, carries the infrared dimension

$$d_{IR}[\mathcal{O}] = n_\partial[\mathcal{O}] + n_g[\mathcal{O}] + \frac{1}{2}n_f[\mathcal{O}], \quad (1)$$

where $n_\partial[\mathcal{O}]$ is the number of derivatives entering the operator \mathcal{O} , $n_g[\mathcal{O}]$ the number of gauge coupling constants and $n_f[\mathcal{O}]$ the number of fermion fields. This infrared counting rule provides the basis for the ordering of diagrams. Provided fermion masses can consistently be counted as $\mathcal{O}(p^1)$ or higher (more on this later), one finds that the degree of suppression $d_{IR}[\Gamma]$ of the diagram Γ is given by⁴

$$d_{IR}[\Gamma] = 2 + 2L + \sum_{v=1}^V (d_{IR}[\mathcal{O}_v] - 2), \quad (2)$$

where L is the number of loops and $d_{IR}[\mathcal{O}_v]$ is the infrared dimension (1) of the vertex \mathcal{O}_v , and the vertices are numbered $v = 1, \dots, V$. The expansion in powers of d_{IR} is also an expansion in loops L : it makes sense at least at the formal level provided *all* operators invariant under the symmetry have $d_{IR}[\mathcal{O}] \geq 2$.

2.2 Unwanted operators at leading order

Among all $SU(2)_L \times U(1)_Y$ -invariant operators of lowest order in the Higgs-less theory, one finds (at leading order) operators which have no equivalent in the renormalizable lagrangian of the SM. In fact, the corresponding operators can be built using SM fields, but they would have mass-dimension six. In the absence of the Higgs particle, this suppression no longer holds.

Using a left-handed lepton doublet ℓ_L , one can translate Weinberg's⁵ LNV $SU(2)_L \times U(1)_Y$ -invariant operator to the Higgs-less case, obtaining $\Lambda \bar{\ell}_L \Sigma \tau^+ \Sigma^\dagger (\ell_L)^c = \mathcal{O}(p^1)$. According to the power-counting rules given above, this operator appears at $\mathcal{O}(p^1)$, without any suppression factor. Other operators that can be written at $\mathcal{O}(p^1)$ yield Dirac masses to fermions $\Lambda \bar{\chi}_L \Sigma (\tau^3) \chi_R$. Such operators have chiral dimension less than two, and therefore endanger the internal consistency of the expansion procedure, as mentioned in Section 2.1.

At $\mathcal{O}(p^2)$, one finds other “unwanted” operators involving fermions. Non-universal couplings^{6,7} to massive vector bosons appear at $\mathcal{O}(p^2)$: $i\overline{\chi_L}\gamma^\mu(\Sigma D_\mu\Sigma^\dagger)\chi_L$ and $i\overline{\chi_R}\gamma^\mu(\Sigma^\dagger D_\mu\Sigma)\chi_R$. In addition, this introduces couplings of the right-handed fermions to the W^\pm . Both types of operators would also be a new source of flavor-changing currents, requiring a redefinition of the CKM matrix, which would not be unitary anymore. These operators would also introduce flavor-changing neutral currents (FCNCs) at this level.

At $\mathcal{O}(p^2)$, we find two more operators, giving tree-level contributions to the $S^{8,9}$ and T^{10} parameters. These operators are $SU(2)_L \times U(1)_Y$ -invariant, but break custodial symmetry.

3 Higgs-less LEET based on S_{nat}

3.1 The symmetry S_{nat}

We assume a larger hidden symmetry $S_{\text{nat}} \supset SU(2)_L \times U(1)_Y$: this allows for an expansion procedure consistent with the principles of a LEET, and in which the unwanted operators of the previous Section are relegated to higher orders. In the minimal version, and before the constraints are applied, the lagrangian of the theory at $\mathcal{O}(p^2)$ consists of two decoupled sectors: a) the symmetry-breaking sector containing three GBs together with six connections of the spontaneously-broken $SU(2)_{\Gamma_L} \times SU(2)_{\Gamma_R}$ symmetry *and* b) an unbroken $SU(2)_{G_L} \times SU(2)_{G_R} \times U(1)_{B-L}$ gauge theory with the $L \leftrightarrow R$ symmetric coupling of local left and right isospin to chiral fermion doublets. The symmetry group S_{nat} is thus

$$S_{\text{nat}} = [SU(2) \times SU(2)]^2 \times U(1)_{B-L}. \quad (3)$$

We assume an underlying theory responsible for the spontaneous symmetry breaking of $SU(2)_{\Gamma_L} \times SU(2)_{\Gamma_R} \subset S_{\text{nat}}$ down to its vector subgroup. This produces a triplet of GBs Σ transforming according to^b

$$\Sigma \mapsto \Gamma_L \Sigma \Gamma_R^\dagger, \quad (4)$$

where $\Gamma_L \in SU(2)_{\Gamma_L}$ and $\Gamma_R \in SU(2)_{\Gamma_R}$, and the corresponding connections are denoted by $\Gamma_{L\mu}, \Gamma_{R\mu}$. So much for the composite sector of the theory.

On the other hand, for the elementary sector of the theory, the elementary fermion doublets $\chi_{L,R}$ transform as

$$\chi_{L,R} \mapsto G_{L,R} e^{-i\frac{B-L}{2}\alpha} \chi_{L,R}. \quad (5)$$

The gauge fields of the elementary group $SU(2)_{G_L} \times SU(2)_{G_R} \times U(1)_{B-L}$ are denoted by $g_L G_{L\mu}, g_R G_{R\mu}, g_B G_{B\mu}$, and we thus collect all $\mathcal{O}(p^2)$ terms as

$$\begin{aligned} \mathcal{L}(p^2) &= \frac{f^2}{4} \left\langle D_\mu \Sigma^\dagger D^\mu \Sigma \right\rangle + i\overline{\chi_L} \gamma^\mu D_\mu \chi_L + i\overline{\chi_R} \gamma^\mu D_\mu \chi_R \\ &\quad - \frac{1}{2} \left\langle G_{L\mu\nu} G_L^{\mu\nu} + G_{R\mu\nu} G_R^{\mu\nu} \right\rangle - \frac{1}{4} G_{B\mu\nu} G_B^{\mu\nu}. \end{aligned} \quad (6)$$

We see that the symmetry S_{nat} eliminates all the unwanted couplings discussed in Section 2.2 at the leading chiral order $\mathcal{O}(p^2)$ described by the lagrangian (6). On the other hand, it seems at first sight that S_{nat} is too large: the lagrangian (6) contains thirteen gauge connections $g_L G_{L\mu}, g_R G_{R\mu}, g_B G_{B\mu}, \Gamma_{L\mu}, \Gamma_{R\mu}$, as compared to four in the SM. Due to the GB term in (6) (first term in the right-hand side), the three combinations $\Gamma_{R\mu}^a - \Gamma_{L\mu}^a$ acquire a mass term by the Higgs mechanism, whereas all ten remaining vector fields as well as fermions remain massless. Furthermore, the lagrangian (6) does not contain any coupling that would transmit the symmetry breaking from the composite sector $(\Sigma, \Gamma_{L\mu}, \Gamma_{R\mu})$ to the elementary sector $(G_{L\mu}, G_{R\mu}, G_{B\mu}, \chi)$. We now remedy this by introducing constraints that reduce the space of gauge connections.

^bThe notation Σ is the same as in the previous Section, but the transformation properties are different.

3.2 Reduction $S_{\text{nat}} \rightarrow S_{\text{red}}$ via constraints

We want to identify $\Gamma_{L\mu}$ to $g_L G_{L\mu}$, up to a gauge transformation $\Omega_L \in \text{SU}(2)$, i.e.

$$\Gamma_{L\mu} = \Omega_L(x) g_L G_{L\mu} \Omega_L^{-1}(x) + i\Omega_L(x) \partial_\mu \Omega_L^{-1}(x). \quad (7)$$

This will reduce the group $\text{SU}(2)_{G_L} \times \text{SU}(2)_{\Gamma_L}$ to its vector subgroup, which will be recognized as the $\text{SU}(2)_L$ of the SM. Requiring the invariance of the constraint (7) with respect to the whole symmetry $\text{SU}(2)_{G_L} \times \text{SU}(2)_{\Gamma_L}$ amounts to promoting the gauge function Ω_L to a field that transforms according to

$$\Omega_L \mapsto \Gamma_L \Omega_L G_L^\dagger. \quad (8)$$

The field Ω_L is *not* a GB, but rather a non-propagating *spurion*: this follows from the constraint (7) itself, since the latter can be equivalently rewritten as

$$D_\mu \Omega_L \equiv \partial_\mu \Omega_L - i\Gamma_{L\mu} \Omega_L + ig_L \Omega_L G_{L\mu} = 0. \quad (9)$$

The spurion has no dynamics, since no kinetic term can be written down for it. A similar procedure can be performed in the right-handed sector: it is complicated by the selection of the $\text{U}(1)$ subgroups³, so we do not describe it here.

The problem of unwanted terms reappears as long as the spurion \mathcal{X} is restricted to be unitary: one can still construct other S_{nat} -invariants that are $\mathcal{O}(p^2)$ but are not contained in (6). These additional terms are exactly all the unwanted terms of Section 2.2, which still have to be suppressed. This can be achieved by adding a new ingredient: we now admit multiplication of the unitary spurions by constants which are (technically) naturally small. This is implemented via the requirement that only the object \mathcal{X}

$$\mathcal{X} \equiv \xi \Omega_L, \quad (10)$$

may be inserted in the operators of Section 2.2 in order to make them invariant under S_{nat} . The order of magnitude of ξ should later be estimated from experiments, but ξ will be considered as an expansion parameter. We now require, by analogy with (9)

$$D_\mu \mathcal{X} \equiv \partial_\mu \mathcal{X} - i\Gamma_{L\mu} \mathcal{X} + ig_L \mathcal{X} G_{L\mu} = 0. \quad (11)$$

This implies the existence of a “standard gauge”, specified by $\Omega_L = 1$, in which the connections are equal, i.e. $\Gamma_{L\mu} \stackrel{\text{s.g.}}{=} g_L G_{L\mu}$: in this gauge the spurion \mathcal{X} reduces to one constant ξ .

For the right-handed sector, one has to define two spurions \mathcal{Y} and \mathcal{Z} . One may again define the standard gauge, in which the spurions reduce to two constants η and ζ respectively, while the gauge connections are identified appropriately according to $\Gamma_{R\mu}^{1,2} \stackrel{\text{s.g.}}{=} g_R G_{R\mu}^{1,2} \stackrel{\text{s.g.}}{=} 0$ and $\Gamma_{R\mu}^3 \stackrel{\text{s.g.}}{=} g_R G_{R\mu}^3 \stackrel{\text{s.g.}}{=} g_B G_{B\mu}^3$.

4 Consequences of the formalism

4.1 Construction of the LEET

The next step in the formulation of the LEET is the construction of the effective lagrangian: one writes down all terms invariant under S_{nat} that can be constructed out of the GBs Σ , the connections $\Gamma_{L\mu}, \Gamma_{R\mu}$, the gauge fields $G_{L\mu}, G_{R\mu}, G_{B\mu}$, the spurions $\mathcal{X}, \mathcal{Y}, \mathcal{Z}$, and fermions. The operators should be ordered according to their chiral power-counting and to the powers of spurions involved. To exhibit the physical content of each operator, one then injects the solution of the constraints in the standard gauge. This yields a lagrangian depending on the fermions

and on the S_{red} gauge fields, which should be used as dynamical variables to compute loops. In addition, this lagrangian depends on the three constants ξ , η and ζ . At the leading order $\mathcal{O}(p^2)$ without explicit powers of spurions, the lagrangian describes exactly the SM couplings but without the Higgs boson, and with all fermions left massless. The origin of fermion masses, which will come with explicit powers of spurions, thus appears different from that of vector bosons. Other terms involving explicit powers of spurions will also bring other interactions³.

4.2 Dirac masses

To be invariant under S_{nat} , Dirac masses require one insertion of the spurion \mathcal{X} and one of the spurion \mathcal{Y} (\mathcal{Y} and $\mathcal{Y}_c \equiv \tau^2 \mathcal{Y}^* \tau^2$ can be conveniently decomposed as linear combinations of \mathcal{Y}_u and \mathcal{Y}_d). Hence, the leading quark mass term in the lagrangian is of order $\mathcal{O}(p^1 \xi^1 \eta^1)$ and reads

$$\mathcal{L}(p^1 \xi^1 \eta^1)_{\text{quarks}} = -\Lambda_{\text{quarks}} \left(\overline{q_L} \mathcal{X}^\dagger \Sigma \mathcal{Y}_u q_R + \overline{q_L} \mathcal{X}^\dagger \Sigma \mathcal{Y}_d q_R \right) + \text{h.c.} \quad (12)$$

The Dirac mass terms for leptons can be written in full analogy with the quark mass term (12), yielding neutrino Dirac mass terms of comparable magnitude.

As mentioned in Section 2.1, consistency of the low-energy power counting for a fermion propagator inside loops requires fermion masses to count as $\mathcal{O}(p^1)$ or smaller. This is possible here, thanks to the occurrence of spurions in the fermion mass terms: this suggests a relation between spurion and momentum expansion, specified by the counting rule

$$\xi \eta = \frac{m_t}{\Lambda_{\text{quarks}}} = \mathcal{O}(p^1). \quad (13)$$

4.3 Neutrino Dirac and Majorana masses

Making use of the spurion \mathcal{Z} , one can construct $\Delta L = 2$ operators that are S_{nat} -invariant. The spurion \mathcal{Z} responsible for the selection of the $U(1)_Y$ subgroup, controls —via the parameter ζ — the strength of these. This leads us to the assume $\zeta \ll \xi, \eta \ll 1$.

Majorana masses of left- and right-handed neutrinos can thus be suppressed in the present LEET, since they involve a coefficient ζ^2 . To obtain left-handed neutrinos lighter than the charged fermions, one can forbid the neutrino Dirac masses mentioned after equation (12): this is done imposing a \mathbb{Z}_2 symmetry which, in the standard gauge simply reduces to $\nu_R \mapsto -\nu_R$.

At this point, it is worth stressing that the *number* of spurions to be introduced is entirely fixed once we have identified the higher S_{nat} symmetry, and once we ask to recover the electroweak group S_{red} by imposing constraints. In the right-handed sector, there would a priori be various possibilities for the introduction of the expansion parameters: the physical requirement that $B - L$ breaking effects be small leaves us with three inequivalent possibilities for the $(B - L)$ -breaking building block \mathcal{Z} . This distinction in turn implies different estimates for the ν_R masses, and therefore different cosmological consequences³.

4.4 Comparing vertex and oblique corrections

The vertex corrections as written in Section 2.2 would not be invariant under S_{nat} . On the other hand, inserting appropriate powers of spurions, we find the following S_{nat} -invariants, of order respectively $\mathcal{O}(p^2 \xi^2)$ and $\mathcal{O}(p^2 \eta^2)$

$$i\overline{\chi_L} \gamma^\mu \mathcal{X}^\dagger (\Sigma D_\mu \Sigma^\dagger) \mathcal{X} \chi_L, \quad i\overline{\chi_R} \gamma^\mu \mathcal{Y}_{u,d}^\dagger (\Sigma^\dagger (D_\mu \Sigma) \mathcal{Y}_{u,d}) \chi_R. \quad (14)$$

This suggests that, in the Higgs-less LEET, certain vertex corrections could be more important than oblique ones, which involve more powers of momenta or spurions³. A parametrization of vertex corrections at NLO, using some simple assumptions about the flavor structure is presented in³: this opens the possibility of looking for non-oblique deviations of the SM.

5 Conclusions

We have constructed a systematic LEET formalism for EWSB without a physical Higgs, which can be renormalized order by order in a momentum expansion. The leading order does not display deviations from the SM. This requires a hidden symmetry $S_{\text{nat}} \supset \text{SU}(2)_L \times \text{U}(1)_Y = S_{\text{red}}$, reduced to the electroweak group S_{red} via constraints. Implementing the constraints in a covariant manner requires spurions: these can be used to introduce the small expansion parameters describing effects beyond the SM.

The spurions live in the coset space $S_{\text{nat}}/S_{\text{red}}$. The constraint implies that spurions do not propagate and do not generate mass terms for vector fields either, in contrast to GBs. There exists a “standard gauge” in which spurions reduce to a set of constants. In the actual case of the group S_{nat} , the spurions reduce to three constants, denoted ξ, η and ζ . This reflects the structure of the coset space which, in this case, is a product of three $\text{SU}(2)$ groups.

The S_{nat} -invariant constraints eliminate the nine redundant fields, reduce the *linear symmetry* S_{nat} to its electroweak subgroup S_{red} , and induce couplings between the symmetry-breaking and gauge/fermion sectors. The W^\pm and Z^0 become massive, whereas all fermions remain massless. In this way one recovers all Higgs-less vertices of the SM. The main effect of S_{nat} is the elimination of all non-standard $\mathcal{O}(p^2)$ vertices.

The expansion parameters ξ, η and ζ play a role similar to quark masses in χ PT. The complete LEET invariant under S_{nat} is defined as a double expansion: in powers of momenta and in powers of spurions. The LEET at leading order coincides with the Higgs-less vertices of the SM, used at tree level. Majorana and Dirac mass terms, which could appear at $\mathcal{O}(p^1)$, now involve in addition powers of spurions: within the expansion, these operators can be consistently counted as $\mathcal{O}(p^2)$ or higher. Also, non-standard $\mathcal{O}(p^2)$ vertices reappear as S_{nat} -invariant operators explicitly containing spurions, i.e. suppressed by powers of the parameters ξ, η and ζ in addition to p^2 . In particular, we note that vertex corrections appear before oblique ones (the latter cannot be disentangled from loops).

The consequences of the assumed higher symmetry can be studied systematically using the formalism proposed here for the Higgs-less case: vertex corrections, lepton-number violation as well as the cosmological consequences of light right-handed neutrinos³.

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